On the cohomological derivation of topological Yang-Mills theory

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Abstract

Topological Yang-Mills theory is derived in the framework of Lagrangian BRST cohomology.

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1 Introduction

The cohomological understanding of the BRST symmetry [1]–[9] allowed, apart from proving the existence of the BRST generator for an arbitrary gauge system [9]–[10], a useful investigation of many interesting aspects related to perturbative renormalization problem [11]–[13], anomaly-tracking mechanism [13]–[14], simultaneous study of local and rigid invariances of a given theory [15], as well as to the construction of consistent interactions in gauge theories [16]–[19]. The last topic is probably the most efficient proof of the power of cohomological BRST ideas, reformulating the classical Lagrangian problem of building consistent interactions in gauge theories in terms of precise cohomological classes of the BRST differential, which further offers a systematic search for all possible consistent interactions in the natural background of the deformation theory. Among the models of great

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interest in theoretical physics that have been inferred along the deformation of the master equation we mention Yang-Mills theory [20], the Freedman-Townsend model [21], the Chapline-Manton model [22]. Also, it is important to notice the deformation results connected to Einstein's gravity theory [23], four- and eleven-dimensional supergravity [24], p-forms [25] or chiral forms [26]. However, there remain some important coupled models that have not been recovered so far in the light of the deformation of the master equation, like the topological Yang-Mills theory. This is precisely the main aim of our paper, namely, to infer the four-dimensional topological coupling among Yang-Mills fields via the deformation technique. In view of this, we begin with a certain uncoupled theory in four dimensions and derive its associated BRST symmetry. Consequently, we write down the equations that should be satisfied by the deformed solution to the master equation in terms of the coupling constant, and find their consistent solutions by using the BRST symmetry for the free model. In this manner, we find a complete deformed solution, which is consistent at all orders in the coupling constant. From the analysis of the structure of this solution we observe that the resulting coupled model is nothing but the topological Yang-Mills theory. Thus, the procedure applied in our paper leads to a nice example of simultaneous deformation of the gauge transformations, gauge algebra and reducibility relations of the starting uncoupled system.

2 BRST symmetry for the uncoupled model

Initially, we infer the antifield-BRST symmetry for an uncoupled model, described by the Lagrangian action

$$S_0^L\left[A_\mu^a\right]=-\frac{1}{4}\int d^4x\varepsilon_{\mu\nu\lambda\rho}F_a^{\mu\nu}F^{a\lambda\rho}. \qquad (1)$$

The field strength is defined by $F_a^{\mu\nu}=\partial^\mu A_a^\nu-\partial^\nu A_a^\mu\equiv\partial^{[\mu}A_a^{\nu]}$, while $\varepsilon_{\mu\nu\lambda\rho}$ denotes the completely antisymmetric four-dimensional symbol. Action (1) is invariant under the gauge transformations

$$\delta_\epsilon\Phi^{\alpha_0}=Z^{\alpha_0}_{\ \alpha_1}\epsilon^{\alpha_1}\rightarrow\delta_\epsilon A_\mu^a=\partial_\mu\epsilon^a+\epsilon_\mu^a, \qquad (2)$$

with

$$\Phi^{\alpha_0}\rightarrow A_\mu^a,\ \epsilon^{\alpha_1}\rightarrow\left(\begin{array}{c}\epsilon^b\\ \epsilon_\nu^b\end{array}\right),\ Z^{\alpha_0}_{\ \alpha_1}\rightarrow\left(\delta^a_{\ b}\partial_\mu,\ \delta^a_{\ b}\delta^\nu_{\ \mu}\right), \qquad (3)$$

which are first-stage reducible. Indeed, if we take $\epsilon^a = \theta^a$, $\epsilon^a_\mu = -\partial_\mu \theta^a$, then the gauge transformations (2) vanish identically, $\delta_\epsilon A^a_\mu = 0$. Consequently, the reducibility relations can be written like

$$Z^{\alpha_0}_{\ \ \alpha_1} Z^{\alpha_1}_{\ \ \alpha_2} = 0, \tag{4}$$

with the first-stage reducibility matrix given by

$$Z^{\alpha_1}_{\ \ \alpha_2} \rightarrow \left(\begin{array}{c} \delta^b_{\ c} \\ -\delta^b_{\ c}\partial_\nu \end{array} \right). \tag{5}$$

Accordingly, the solution to the master equation of the uncoupled model is expressed by

$$S = S_0^L \left[A^a_\mu \right] + \int d^4x \left(A^{*\mu}_a \left(\partial_\mu C^a + C^a_\mu \right) + C^*_a \eta^a - C^{*\mu}_a \partial_\mu \eta^a \right), \tag{6}$$

where C^a and C^a_μ stand for the fermionic ghost number one ghosts, and η^a denote the bosonic ghost number two ghosts required by the reducibility. The star variables $A^{*\mu}_a$, C^*_a, $C^{*\mu}_a$ and η^*_a represent the antifields of the corresponding fields/ghosts. The antifields $A^{*\mu}_a$ are fermionic with ghost number minus one, C^*_a and $C^{*\mu}_a$ are bosonic with ghost number minus two, while η^*_a are fermionic and display ghost number minus three. The ghost number is defined in the standard manner like the difference between pure ghost number (pgh) and antighost number (antigh), with

$$\mathrm{pgh}\left(A^a_\mu \right) = \mathrm{pgh}\left(A^{*\mu}_a \right) = \mathrm{pgh}\left(C^*_a \right) = \mathrm{pgh}\left(C^{*\mu}_a \right) = \mathrm{pgh}\left(\eta^*_a \right) = 0, \tag{7}$$

$$\mathrm{pgh}\left(C^a \right) = \mathrm{pgh}\left(C^a_\mu \right) = 1,\ \mathrm{pgh}\left(\eta^a \right) = 2, \tag{8}$$

$$\mathrm{antigh}\left(A^a_\mu \right) = \mathrm{antigh}\left(C^a \right) = \mathrm{antigh}\left(C^a_\mu \right) = \mathrm{antigh}\left(\eta^a \right) = 0, \tag{9}$$

$$\mathrm{antigh}\left(A^{*\mu}_a \right) = 1,\ \mathrm{antigh}\left(C^{*\mu}_a \right) = \mathrm{antigh}\left(C^*_a \right) = 2,\ \mathrm{antigh}\left(\eta^*_a \right) = 3. \tag{10}$$

The antifield BRST differential $s\bullet = (\bullet, S)$ of the free model splits as

$$s = \delta + \gamma, \tag{11}$$

where δ is the Koszul-Tate differential, and γ denotes the longitudinal exterior derivative along the gauge orbits. The symbol $(,)$ signifies the antibracket in the antifield formalism. Consequently, we find that

$$\delta A^a_\mu = 0,\ \gamma A^a_\mu = \partial_\mu C^a + C^a_\mu, \tag{12}$$

$$\delta C^a = 0,\ \gamma C^a = \eta^a, \tag{13}$$

$$\delta C^a_\mu = 0,\ \gamma C^a_\mu = -\partial_\mu \eta^a, \tag{14}$$

$$\delta \eta^a = 0,\ \gamma \eta^a = 0, \tag{15}$$

$$\delta A^{*\mu}_a = 0,\ \gamma A^{*\mu}_a = 0, \tag{16}$$

$$\delta C^*_a = -\partial_\mu A^{*\mu}_a,\ \gamma C^*_a = 0, \tag{17}$$

$$\delta C^{*\mu}_a = A^{*\mu}_a,\ \gamma C^{*\mu}_a = 0, \tag{18}$$

$$\delta \eta^*_a = -\left(C^*_a + \partial_\mu C^{*\mu}_a\right),\ \gamma \eta^*_a = 0. \tag{19}$$

The above formulas will be employed in the next section at the deformation of the solution (6) in a cohomological context.

3 Deformation procedure

A consistent deformation of action (1) and of its gauge invariances defines a deformation of the solution to the master equation that preserves both the master equation and the field/antifield spectra [16]. This means that if

$$S^L_0\left[A^a_\mu\right] + g\int d^4x\alpha_0 + g^2\int d^4x\beta_0 + O\left(g^3\right), \tag{20}$$

is a consistent deformation of action (1), with deformed gauge transformations

$$\bar{\delta}_\epsilon A^a_\mu = \partial_\mu \epsilon^a + \epsilon^a_\mu + g\lambda^a_\mu + O\left(g^2\right), \tag{21}$$

then the deformed solution to the master equation

$$\bar{S} = S + g\int d^4x\alpha + g^2\int d^4x\beta + O\left(g^3\right) = S + gS_1 + g^2S_2 + O\left(g^3\right), \tag{22}$$

should satisfy

$$\left(\bar{S}, \bar{S}\right) = 0, \tag{23}$$

where

$$\alpha = \alpha_0 + A^{*\mu}_a \bar{\lambda}^a_\mu + \text{“more”}. \tag{24}$$

The master equation (23) splits according to the deformation parameter g as

$$s\alpha = \partial_\mu j^\mu, \tag{25}$$

$$s\beta + \frac{1}{2}\omega = \partial_\mu \theta^\mu, \tag{26}$$

$$\vdots$$

for some local j^μ and θ^μ, with

$$(S_1, S_1) = \int d^4 x \omega. \tag{27}$$

We omitted the zeroth order equation in the coupling constant corresponding to the (23) as this is automatically verified. From (25) we read that the first-order non-trivial consistent deformations belong to $H^0(s|d)$, where d is the exterior space-time derivative. In the situation where α is a coboundary modulo d ($\alpha = s\lambda + \partial_\mu \pi^\mu$), the corresponding deformation is trivial (it can be eliminated by a redefinition of the fields).

In order to solve equation (25), we expand α accordingly the antighost number

$$\alpha = \alpha_0 + \alpha_1 + \cdots + \alpha_I, \ \text{antigh}(\alpha_K) = K, \tag{28}$$

where the last term in (28) can be assumed to be annihilated by γ. Since $\text{antigh}(\alpha_I) = I$ and $\text{gh}(\alpha_I) = 0$, it follows that $\text{pgh}(\alpha_I) = I$. Therefore, we can represent α_I under the form

$$\alpha_I = \mu_{a_1 \cdots a_M b_1 \ldots b_N} \eta^{a_1} \cdots \eta^{a_M} \rho^{b_N} \cdots \rho^{b_1}, \tag{29}$$

where N and M are some nonnegative integers satisfying $3N + 2M = I$, and $\mu_{a_1 \cdots a_M b_1 \ldots b_N}$ stand for some γ-invariant functions with $\text{antigh}(\mu_{a_1 \cdots a_N b_1 \ldots b_M}) = 3N + 2M$. In (28), we used the notation

$$\rho^a = f^a_{\ bc} \eta^b C^c, \tag{30}$$

with $f^a_{\ bc}$ some constants that are antisymmetric in the lower indices, $f^a_{\ bc} = -f^a_{\ cb}$. The antisymmetry of $f^a_{\ bc}$ is implied by the γ-invariance of ρ^a. Thus, the general form of $\mu_{a_1 \cdots a_M b_1 \ldots b_N}$ is given by

$$\mu_{a_1 \cdots a_M b_1 \ldots b_N} = (-)^N \eta^*_{b_1} \cdots \eta^*_{b_N} \left(\sum_{k=0}^{M} C^*_{a_1} \cdots C^*_{a_{M-k}} \left(\partial_{\mu_1} C^{*\mu_1}_{a_{M-k+1}} \right) \cdots \left(\partial_{\mu_k} C^{*\mu_k}_{a_M} \right) \right), \tag{31}$$

so

$$\alpha_I = \rho^N \sum_{k=0}^{M} \lambda^{M-k} \sigma^k, \tag{32}$$

where

$$\rho=-\eta_a^*\rho^a,\ \lambda=C_a^*\eta^a,\ \sigma=\left(\partial_\mu C_a^{*\mu}\right)\eta^a. \tag{33}$$

Taking into account the formulas (12)–(19) and the above form of α_I, we obtain

$$\begin{aligned}
&\delta\alpha_I=\gamma\left(\left(\rho^N\sum_{k=0}^{M}k\lambda^{M-k}\sigma^{k-1}-(M-k)\,\lambda^{M-k-1}\sigma^k\right)A_b^{*\mu}C_\mu^b+\right.\\
&\left.N\rho^{N-1}\left(\sum_{k=0}^{M}\lambda^{M-k}\sigma^k\right)\left(\frac{1}{2}C_a^*f^a_{\ bc}C^bC^c+C_a^{*\mu}f^a_{\ bc}\left(C^bC_\mu^c+\eta^bA_\mu^c\right)\right)\right)+\\
&\partial_\mu\left(\left(\rho^N\sum_{k=0}^{M}k\lambda^{M-k}\sigma^{k-1}-(M-k)\,\lambda^{M-k-1}\sigma^k\right)A_b^{*\mu}\eta^b+\right.\\
&\left.N\rho^{N-1}\left(\sum_{k=0}^{M}\lambda^{M-k}\sigma^k\right)C_a^{*\mu}f^a_{\ bc}C^b\eta^c\right)-\\
&A_b^{*\mu}\eta^b\partial_\mu\left(\rho^N\sum_{k=0}^{M}k\lambda^{M-k}\sigma^{k-1}-(M-k)\,\lambda^{M-k-1}\sigma^k\right)-\\
&C_a^{*\mu}f^a_{\ bc}C^b\eta^c\partial_\mu\left(N\rho^{N-1}\sum_{k=0}^{M}\lambda^{M-k}\sigma^k\right).
\end{aligned} \tag{34}$$

On the other hand, equation (25) at antighost number $(I-1)$ takes the form

$$\delta\alpha_I+\gamma\alpha_{I-1}=\partial_\mu m^\mu, \tag{35}$$

for some m^μ. The equations (34) and (35) have to be compatible. This happens if and only if $M=0$ and $N=1$, such that $I=3$. In this way, we inferred that the last term in (28) is given precisely by

$$\alpha_3=\rho=-\eta_a^*f^a_{\ bc}\eta^bC^c. \tag{36}$$

Now, from (34) restricted to $M=0$ and $N=1$, we find

$$\begin{aligned}
&\delta\alpha_3=\gamma\left(\frac{1}{2}C_a^*f^a_{\ bc}C^bC^c+C_a^{*\mu}f^a_{\ bc}\left(C^bC_\mu^c+\eta^bA_\mu^c\right)\right)+\\
&\partial_\mu\left(C_a^{*\mu}f^a_{\ bc}C^b\eta^c\right),
\end{aligned} \tag{37}$$

such that

$$\alpha_2=-\frac{1}{2}C_a^*f^a_{\ bc}C^bC^c-C_a^{*\mu}f^a_{\ bc}\left(C^bC_\mu^c+\eta^bA_\mu^c\right). \tag{38}$$

With α_2 at hand, we determine α_1 as solution to the equation

$$\delta\alpha_2 + \gamma\alpha_1 = \partial_\mu n^\mu. \tag{39}$$

On behalf of (38), we get

$$\delta\alpha_2 = \frac{1}{2}\left(\partial_\mu A_a^{*\mu}\right) f^a_{\ bc} C^b C^c - A_a^{*\mu} f^a_{\ bc}\left(C^b C_\mu^c + \eta^b A_\mu^c\right). \tag{40}$$

Then, the solution to (39) is expressed by

$$\alpha_1 = A_a^{*\mu} f^a_{\ bc} A_\mu^c C^b, \tag{41}$$

which yields to

$$\delta\alpha_2 + \gamma\alpha_1 = \partial_\mu\left(\frac{1}{2} A_a^{*\mu} f^a_{\ bc} C^b C^c\right). \tag{42}$$

By projecting (25) on antighost number one, we deduce the relation

$$\delta\alpha_1 + \gamma\alpha_0 = \partial_\mu v^\mu. \tag{43}$$

In the meantime, as $\delta\alpha_1 = 0$, we arrive at

$$\alpha_0 = \frac{1}{2}\varepsilon_{\mu\nu\lambda\rho} f^a_{\ bc} F_a^{\mu\nu} A^{\lambda b} A^{\rho c}, \tag{44}$$

that further leads to

$$\gamma\alpha_0 = \partial_\mu\left(\varepsilon^{\mu\nu\lambda\rho} f^a_{\ bc}\left(C_{\rho a} A_\nu^b A_\lambda^c + C_a \partial_\rho\left(A_\nu^b A_\lambda^c\right)\right)\right). \tag{45}$$

Putting together the results expressed by (36), (38), (41) and (44), it results that the complete first-order deformation reads as

$$S_1 = \int d^4x \left(\frac{1}{2}\varepsilon_{\mu\nu\lambda\rho} f^a_{\ bc} F_a^{\mu\nu} A^{\lambda b} A^{\rho c} + A_a^{*\mu} f^a_{\ bc} A_\mu^c C^b - \right.$$
$$\left.\frac{1}{2} C_a^* f^a_{\ bc} C^b C^c - C_a^{*\mu} f^a_{\ bc}\left(C^b C_\mu^c + \eta^b A_\mu^c\right) - \eta_a^* f^a_{\ bc} \eta^b C^c\right). \tag{46}$$

Until now we proved the existence of α as solution to the equation (25), which is equivalent with the consistency of the interaction to order g. The interaction is also consistent to order g^2 if and only if equation (26) possesses solution (with respect to β), hence if and only if ω, introduced through (27), is s-exact modulo d. By direct computation, we infer

$$(S_1, S_1) = f^a{}_{[bc} f^c{}_{de]} \int d^4x \left(\eta^*_a C^b C^e \eta^d + C^{*\mu}_a C^b C^d C^e_\mu - \right.$$
$$2C^{*\mu}_a C^d \eta^b A^e_\mu + \frac{1}{3} C^*_a C^b C^d C^e + \left. \varepsilon^{\mu\nu\lambda\rho} A^b_\nu A^e_\rho F^{\mu\nu}_a C^d \right) +$$
$$2f^a{}_{bc} f_{ade} \varepsilon^{\mu\nu\lambda\rho} \int d^4x A^b_\lambda A^c_\rho A^e_\nu \partial_\mu C^d \equiv \int d^4x \omega. \tag{47}$$

From (47) we find that ω is s-exact modulo d if and only if the constants $f^a{}_{bc}$ fulfill the Jacobi identity

$$f^a{}_{[bc} f^c{}_{de]} = 0. \tag{48}$$

Consequently, it follows that

$$\beta = -\frac{1}{4} f^a{}_{bc} f_{ade} \varepsilon^{\mu\nu\lambda\rho} A^b_\mu A^c_\nu A^d_\lambda A^e_\rho, \tag{49}$$

which gives the piece of order g^2 from the deformed solution to the master equation under the form

$$S_2 = -\frac{1}{4} f^a{}_{bc} f_{ade} \int d^4x A^b_\mu A^c_\nu A^d_\lambda A^e_\rho. \tag{50}$$

By direct computation we find that $(S_2, S_2) = 0$, such that the higher-order equations in the deformation parameter are satisfied under the choice

$$S_3 = S_4 = \cdots = 0. \tag{51}$$

By virtue of the above results, we can conclude that the complete solution to the master equation (23) defining our deformation problem is expressed by

$$\bar{S} = \int d^4x \left(-\frac{1}{4} \varepsilon_{\mu\nu\lambda\rho} H^{\mu\nu}_a H^{a\lambda\rho} + A^{*\mu}_a \left((D_\mu)^a{}_b C^b + C^a_\mu \right) - \frac{1}{2} g C^*_a f^a{}_{bc} C^b C^c - \right.$$
$$\left. g C^{*\mu}_a f^a{}_{bc} C^b C^c_\mu + C^*_a \eta^a - C^{*\mu}_a (D_\mu)^a{}_b \eta^b - g \eta^*_a f^a{}_{bc} \eta^b C^c \right), \tag{52}$$

where the notation $(D_\mu)^a{}_b$ stands for the covariant derivative

$$(D_\mu)^a{}_b = \delta^a{}_b \partial_\mu + g f^a{}_{bc} A^c_\mu, \tag{53}$$

while the deformed field strength $H^{a\mu\nu}$ is given by

$$H^{a\mu\nu} = F^{a\mu\nu} - g f^{a}_{\ \ bc} A^{\mu b} A^{\nu c}. \tag{54}$$

Let us analyze now the deformed theory, described by (52). We observe that the antifield-independent piece in (52)

$$\bar{S}_0 = -\frac{1}{4} \int d^4x \varepsilon_{\mu\nu\lambda\rho} H_a^{\mu\nu} H^{a\lambda\rho}, \tag{55}$$

describes nothing but the topological coupling between the vector potentials A_μ^a, known as topological Yang-Mills theory. The structure of the terms linear in the antifields $A_a^{*\mu}$ shows that our procedure deforms also the gauge transformations

$$\bar{\delta}_\epsilon A_\mu^a = \left(D_\mu\right)^a_{\ \ b} \epsilon^b + \epsilon_\mu^a. \tag{56}$$

Moreover, from the terms $-\frac{1}{2} g C_a^* f^a_{\ \ bc} C^b C^c - g C_a^{*\mu} f^a_{\ \ bc} C^b C_\mu^c$ we learn that the resulting gauge algebra is deformed, while the presence of $C_a^* \eta^a - C_a^{*\mu} \left(D_\mu\right)^a_{\ \ b} \eta^b$ indicates that the reducibility functions are also deformed

$$\bar{Z}^{\alpha_1}_{\ \ \alpha_2} = \left(\delta^a_{\ \ b}, -\left(D_\mu\right)^a_{\ \ b}\right). \tag{57}$$

In conclusion, the deformation problem studied in this paper generates the deformations of the gauge transformations, gauge algebra and reducibility relations with respect to the starting uncoupled model.

4 Conclusion

To conclude with, in this paper we have investigated the consistent interaction that can be introduced among a set of topologically coupled free vector fields. Starting with the BRST differential for the free theory, $s = \delta + \gamma$, we initially compute the consistent first-order deformation with the help of some cohomological arguments related to the free model. Next, we prove that the deformation is also second-order consistent and, moreover, matches the higher-order deformation equations. As a result, we are led precisely to the topological Yang-Mills theory, that implies the deformation of the gauge transformations, their algebra and of the accompanying reducibility relations.

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